

# Coherent Destruction of Tunneling in a Two Electron Double Quantum Dot: Interplay of Coulomb interaction, spin-orbit interaction and AC field

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We study a double quantum dot system with two interacting electrons in the presence of a time-dependent periodic (AC) electric field and spin-orbit interaction. We focus on the phenomenon of Coherent Destruction of Tunneling (CDT) for an initially localized state. Because of the periodicity introduced by the AC field we use Floquet theory to find quasi energies with their crossing and anti-crossing points corresponding to CDT. We observe that the AC field rescales the spin orbit and hopping amplitudes in terms of Bessel functions. The zeros of the Bessel functions are of the form of a ratio of AC field strength to its frequency and quasi energies at these points form anti-crossings in our case. We first prepare the system in a triplet state and study the evolution of its probability in the presence of spin-orbit interaction alone. We observe an oscillatory behavior which indicates spin flip transitions. However when the AC field is turned on the probability oscillations are highly suppressed at anti-crossing points leading to localization of initial state and the state retains its memory of spin even in the presence of spin flip interaction.

## I. INTRODUCTION

Semiconductor quantum dots are small conductive regions in a semiconductor with tunable number of electrons. Their shape and size can be varied by varying the gate voltage. Quantum dots are potential candidates for building blocks of a quantum computer [1, 2]. For this to work, it is necessary to couple the dots coherently and keep this coherence for longer times. In this regard, spin dephasing time is of special interest. Studies of spin dephasing time have been carried out in the presence of a magnetic field in [3, 4] on single quantum dots (QDs). Double Quantum Dots (DQDs), where tunneling between the dots is an important additional parameter, introduce new energy scales which open avenues of rich physics with possible applications in spintronics. In DQDs, the spins can be manipulated in many ways such as by electric fields in the presence of a constant magnetic field, taking advantage of both spin-orbit and nuclear hyperfine interactions [5–9]. In these systems, spin manipulation in the presence of AC electric field has also been investigated by considering the periodic motion of the electron leading to electron dipole spin resonance (ESDR) [10]. Prior to our work, spin manipulation of single electron in a QD has been studied in the presence of radiation in THz range [11]; the authors have investigated the electronic Density of States (DOS) in the presence of both spin-orbit interaction and electromagnetic radiation. Other studies consider the effect of an external oscillating electric field on Rabi spin oscillation of a single electron in a DQD [12]. This work shows that Rabi spin oscillation can be slowed at high electric field. They have also shown a decrease of displacement of electron which results in slow driving of spin with increase in electric field. This shows strong dependence of efficiency and time scale of spin manipulation on the external field.

In this work, we investigate spin flip processes due to spin-orbit interaction in a DQD system with two interacting electrons in the presence of external AC electric field. Spin-orbit interaction is the main source of spin flip for the three- and two-dimensional electron states in GaAs-type crystals [13, 14]. The unit cell in this material has no inversion center which gives rise to an effective spin-orbit coupling in the electron spectrum. We show that Coherent Destruction of Tunneling (CDT) occurs in our system. The observation of CDT means suppression of spin flip processes which can be attributed to the AC field localizing the initial state. In this regard our results present a different and more general view of spin dynamics compared to [12] where only a single electron is considered and spin oscillation suppression is observed.

Because of the periodicity introduced by the externally applied field Floquet formalism can be used to investigate the evolution of the system. The quasi-energies are calculated by finding the one period propagator and the dynamics of the system are determined for time scale much larger than the driving period. On this time scale the behavior of the system is completely described by the quasi-energies. In the presence of electric field alone, the dynamics of the single electron in a DQD have been rigorously studied with the successful use of this technique [15]. The effect of crossing and anti-crossing of quasi-energies on the dynamics of electrons was studied in [16] and [17, 18]. However, any such scheme based on Hubbard model separates the singlet and triplet subspaces and consequently the dynamics does not show spin non-conserving transitions. The presence of spin-orbit interaction allows for mixing of singlet and triplet states which are completely decoupled otherwise. We choose spin-orbit field direction such that  $z$  axis in spin space lies parallel to spin-orbit field,  $\Omega$ ; in which case two of the triplet states are decoupled from the remaining spin states [19, 20]. We show that the effect of this field is to rescale tunneling (hopping) and spin-orbit interaction

parameters in terms of Bessel functions. This rescaling depends on initial values of these parameters and ratio of electric field strength to its frequency. This leads to CDT of the initial state when the ratio of field strength and frequency is tuned to zeros of the Bessel function.

We show that the initial state prepared as a spin triplet can transition into singlet subspace due to spin-orbit coupling. Our main focus is to see how this transition can be suppressed by the application of an external electric field. This we relate to the behavior of the initial state at the points near the anti crossings.

This paper is organized as follows. In Sec. II, we develop the model for two interacting electrons confined in a DQD and define a suitable basis for the interacting electrons. In Sec. III, we introduce spin-orbit interaction in the second quantized formalism and define an appropriate geometry for our model. In the next section we give a brief introduction to Floquet theory. Using this formalism we then define Floquet states and quasi-energies in Sec. V. We investigate crossings and anti crossings of quasi-energies and study CDT corresponding to these points. Parameters that tune CDT leading to localization of electrons in our system are discussed. For this purpose, we prepare the system in an initial state and study its time evolution. Finally, we present our results in Sec. VI.

## II. MODEL

We consider two electrons confined in a *DQD* system, Fig.1. The dots lie in the plane  $(\xi, z)$  with tunneling that can occur in the  $\xi$  direction. The spin-orbit field  $\Omega$  points along the  $z$  axis which also defines the spin quantization axis. The dots are detuned by externally applied voltage. Fig.2 shows the effect of detuning on the quantum dot levels. We assume that the electrons are confined near the minima of a double well potential  $V_{QD}$  created by electrical gating of the system. Since there are two electrons in the system, the most suitable basis is  $(n, m)$  where  $n$  and  $m$  denote the number of electrons in the left and right dot respectively. In the second quantized notation, including spin, the states can be written as

$$|(2, 0)S\rangle = c_{L\uparrow}^\dagger c_{L\downarrow}^\dagger |0\rangle \quad (1)$$

$$|(0, 2)S\rangle = c_{R\uparrow}^\dagger c_{R\downarrow}^\dagger |0\rangle \quad (2)$$

$$|(1, 1)S\rangle = \frac{1}{\sqrt{2}}(c_{L\uparrow}^\dagger c_{R\downarrow}^\dagger - c_{L\downarrow}^\dagger c_{R\uparrow}^\dagger) |0\rangle \quad (3)$$

$$|T_+\rangle = c_{L\uparrow}^\dagger c_{R\uparrow}^\dagger |0\rangle \quad (4)$$

$$|T_0\rangle = \frac{1}{\sqrt{2}}(c_{L\uparrow}^\dagger c_{R\downarrow}^\dagger + c_{L\downarrow}^\dagger c_{R\uparrow}^\dagger) |0\rangle \quad (5)$$

$$|T_-\rangle = c_{L\downarrow}^\dagger c_{R\downarrow}^\dagger |0\rangle \quad (6)$$

where  $S$  and  $T$  donate the singlet and triplet spin states.  $c_i^\dagger(c_i)$  are the usual creation (annihilation) operators. The orbital parts of the wavefunction for  $|(1, 1)S\rangle$  and  $|T_{0,\pm}\rangle$  are given by

$$\Psi_\pm^s(r_1, r_2) = \frac{1}{\sqrt{2}}[\Phi_L(r_2)\Phi_R(r_2) \pm \Phi_R(r_1)\Phi_L(r_2)] \quad (7)$$

The orbital parts of the other two singlet states  $|(0, 2)S\rangle$  and  $|(2, 0)S\rangle$ , which represent double occupation on right and left dot respectively, are given by

$$\Psi_{L,R}^d(r_1, r_2) = \Phi_{L,R}(r_1)\Phi_{L,R}(r_2) \quad (8)$$

The orbital functions  $\Psi_+^s(r_1, r_2)$  and  $\Psi_{L,R}^d$  are symmetric under exchange of particles while  $\Psi_-^s(r_1, r_2)$  is antisymmetric under exchange.  $\Phi_L$  and  $\Phi_R$  are Wannier orbitals centered on the left and right dot respectively [21]. In our model the energies of the interacting electrons are detuned by an amount  $2\varepsilon$ . This detuning potential  $V_{bias}$  can be established by the electrostatic gates which also create  $V_{QD}$ . The detuning  $\varepsilon$  represents an energy difference for an electron occupying the left or the right dot. For the symmetric case the electrostatic gates voltages are adjusted such that an electron would have same energy in either of the two wells, see Fig.2. The electrons in the DQD are coupled through Coulomb interaction  $U$  which is the cost of double occupation in states  $|(2, 0)S\rangle$  and  $|(0, 2)S\rangle$ . Hopping between the dots is given by  $W$  and repulsion of singlet(triplet) state with single electron in each dot is given by  $V_+(V_-)$ . In addition, we subject the DQD system to an AC electric field in the presence of spin orbit interaction. The oscillating electric field can be established by the electrical gates to which an AC signal is supplied from external source. The AC field,

confinement potential  $V_{QD}$ , Coulomb interaction, the detuning  $\varepsilon$  and hopping constitute the spin independent part of the Hamiltonian:

$$H_0 = (-1)^i \sum_{i,\sigma} (\varepsilon + E(t)) c_{i,\sigma}^\dagger c_{i,\sigma} + W \sum_{\sigma} (c_{L,\sigma}^\dagger c_{R,\sigma} + c_{R,\sigma}^\dagger c_{L,\sigma}) + U \sum_{i,\sigma} n_{i,\uparrow} n_{i,\downarrow} \quad (9)$$

where we have introduced the spin index  $\sigma$ . The summation on  $i$  runs over left and right dots, i.e 1, 2 respectively.  $E(t) = V \cos(\omega t)$  represents the time dependent electric field with  $V$  and  $\omega$  represents the strength and frequency of the oscillating field respectively.

### III. THE SPIN-ORBIT INTERACTION

In the basis defined in Eqns.(1)–(6), the Hamiltonian, Eqn.(9), can be expressed in the singlet and triplet blocks:

$$H_0 = \begin{pmatrix} H_{SS} & 0 \\ 0 & H_{TT} \end{pmatrix} \quad (10)$$

where the singlet Hamiltonian  $H_{SS}$  in the basis  $|(2, 0)S\rangle$ ,  $|(0, 2)S\rangle$  and  $|(1, 1)S\rangle$  is

$$H_{SS} = \begin{pmatrix} U - \varepsilon - V \cos \omega t & \sqrt{2}W & 0 \\ 0 & U + \varepsilon + V \cos \omega t & \sqrt{2}W \\ \sqrt{2}W & \sqrt{2}W & V_+ \end{pmatrix} \quad (11)$$

and  $H_{TT}$  represents the diagonal triplet subspace Hamiltonian.

The transitions between singlet and triplet states can be mediated by the spin-orbit interaction which can be expressed in second quantized form[19]. The spin-orbit interaction within the space of low energy single electron orbitals in DQD is expressed in terms of the spin-orbit field  $\mathbf{\Omega}$ ,

$$H_{SO} = i \mathbf{\Omega} \cdot \sum_{\alpha,\beta=\uparrow\downarrow} (c_{\alpha,L}^\dagger \sigma^{\alpha\beta} c_{\beta,R} - H.C) \quad (12)$$

where  $\mathbf{\Omega}$  is the spin orbit field and is given by

$$i \mathbf{\Omega} = \langle \Phi_L | P_\xi | \Phi_R \rangle a_\Omega \quad (13)$$

It depends on the orientation of the dots with respect to crystallographic axes through the vector  $a_\Omega$ [22, 23].  $\sigma$  represents the Pauli spin matrix. For a 2DEG in the (001) plane,  $a_\Omega$  is given by

$$a_\Omega = (\beta - \alpha) \cos \theta \mathbf{e}_{[\bar{1}10]} + (\beta + \alpha) \sin \theta \mathbf{e}_{[110]} \quad (14)$$

where the angle between  $\mathbf{e}_\xi$  direction and the [110] crystallographic axis is denoted by  $\theta$ . The matrix element of  $P_\xi$ , the momentum component along the  $\xi$  direction that connects the two dots, is taken between the corresponding Wannier orbitals, and it depends on the envelope wave function and the double dot binding potential [9]. The spin-orbit interaction causes transition between singlet and triplet states taking the single occupation of both dots to double occupation of either of the two dots. The explicit expression for  $\mathbf{\Omega}$  is [21],

$$\mathbf{\Omega} = \frac{4W}{3} \frac{l}{\Lambda_{SO}} \frac{a_\Omega}{|a_\Omega|} \quad (15)$$

where  $l$  is inter-dot distance,  $\Lambda_{SO}$  is the spin-orbit length. Inter-dot distance and hopping amplitude are DQD geometry dependent parameters where as the spin-orbit length  $\Lambda_{SO}$  is determined by the material properties and by the orientation of the DQD with respect to the crystallographic axis. Introducing spin-orbit interaction in the Hamiltonian Eqn.(9), we obtain the complete Hamiltonian for our system

$$\begin{aligned} H(t) &= H_0 + H_{SO} \\ &= (-1)^i \sum_{i,\sigma} (\varepsilon + E(t)) c_{i,\sigma}^\dagger c_{i,\sigma} + W \sum_{\sigma} (c_{L,\sigma}^\dagger c_{R,\sigma} + c_{R,\sigma}^\dagger c_{L,\sigma}) + U \sum_{i,\sigma} n_{i,\uparrow} n_{i,\downarrow} \\ &\quad + i \mathbf{\Omega} \cdot \sum_{\alpha,\beta=\uparrow\downarrow} (c_{\alpha,L}^\dagger \sigma^{\alpha\beta} c_{\beta,R} - h.c). \end{aligned} \quad (16)$$

With the z-axis taken along  $\mathbf{\Omega}$  the Hamiltonian given in Eqn.(16) becomes

$$H(t) = \begin{pmatrix} U - \varepsilon - V \cos \omega t & 0 & \sqrt{2}W & 0 & -\iota\sqrt{2}\Omega & 0 \\ 0 & U + \varepsilon + V \cos \omega t & \sqrt{2}W & 0 & -\iota\sqrt{2}\Omega & 0 \\ \sqrt{2}W & \sqrt{2}W & V_+ & 0 & 0 & 0 \\ 0 & 0 & 0 & \varepsilon_{T+} & 0 & 0 \\ \iota\sqrt{2}\Omega & \iota\sqrt{2}\Omega & 0 & 0 & V_- & 0 \\ 0 & 0 & 0 & 0 & 0 & \varepsilon_{T-} \end{pmatrix} \quad (17)$$

where  $\varepsilon_{T+}$  and  $\varepsilon_{T-}$  refer to energies of triplet states  $|T_+\rangle$  and  $|T_-\rangle$  which are degenerate in energy and do not couple to singlet subspace and  $|T_0\rangle$  state, due to our choice of the geometry of the dots and spin-orbit field  $\mathbf{\Omega}$  axis. In case of a choice of other directions of  $\mathbf{\Omega}$ , the singlet subspace would couple with  $|T_+\rangle$ ,  $|T_-\rangle$  also. In the block diagonal form we have

$$\begin{pmatrix} H_1(t) & 0 \\ 0 & H_2 \end{pmatrix} \quad (18)$$

The matrix  $H_1(t)$  is

$$\begin{pmatrix} U - \varepsilon - V \cos \omega t & 0 & \sqrt{2}W & -i\sqrt{2}\Omega \\ 0 & U + \varepsilon + V \cos \omega t & \sqrt{2}W & -i\sqrt{2}\Omega \\ \sqrt{2}W & \sqrt{2}W & V_+ & 0 \\ i\sqrt{2}\Omega & i\sqrt{2}\Omega & 0 & V_- \end{pmatrix} \quad (19)$$

$H_1(t)$  consists of the singlet states  $|(2,0)S\rangle, |(0,2)S\rangle, |(1,1)S\rangle$  and  $|T_0\rangle$  whereas  $H_2$  is a  $2 \times 2$  matrix in the basis of triplet states  $|T_+\rangle$  and  $|T_-\rangle$ .

#### IV. FLOQUET FORMALISM

The Hamiltonian in Eqn.(19) is periodic in time:  $H_1(t+T) = H_1(t)$  with  $T$  being the period of the driving field. Floquet formalism can be used to express solutions of the time-dependent Schrodinger equation as  $\Psi_\alpha(x, t) = \exp(-i\varepsilon_\alpha t)\phi_\alpha(x, t)$ , where  $\phi_\alpha(x, t)$  is a function with the same periodicity as  $H_1(t)$  and is called a Floquet state and  $\varepsilon_\alpha$  is termed the quasi-energy. The Floquet states and their quasi-energies can be obtained from the eigenvalue equation:

$$H_F \phi_\alpha(x, t) = \varepsilon_\alpha \phi_\alpha(x, t) \quad (20)$$

where  $H_F = (H_I(t) - i\frac{\partial}{\partial t})$  is called Floquet Hamiltonian. It is evident that the Floquet states are periodic in time with the same period as the Hamiltonian. The periodicity of  $\phi_\alpha(x, t)$  allows a series of Floquet eigenvalues,  $\varepsilon_\alpha + n\hbar\omega$ , with corresponding eigenfunctions  $\phi_{n\alpha}(x, t) = e^{(in\omega t)}\phi_\alpha(x, t)$ , where  $n$  is an integer. However, the physical state  $\Psi_\alpha(x, t)$  is unchanged; the Floquet eigenvalues associated with distinct physical states are defined only modulo  $\hbar\omega$ . The quasi-energies can be obtained by numerical diagonalization of the unitary evolution operator  $U(T, 0)$  for one period of the field. The eigenvalues of  $U(T, 0)$  are then related to quasi-energies as  $\lambda_\alpha = \exp[-i\varepsilon_\alpha T]$ . The time periodicity of the Floquet states allows us to study the dynamics of the system for time scales larger than the period of the driving field effectively in terms of quasi-energies. As the quasi-energies approach degeneracy, the dynamics of the system are frozen, producing CDT [18].

#### V. CALCULATION AND RESULTS

To find the approximate solution of Eqn.(19) we follow the perturbation scheme of [24]. We divide the Hamiltonian Eqn.(18) into two parts:  $H_I(t)$  which contains the electric field, detuning and Coulomb terms and  $H_t$  which contains the tunneling component and spin-orbit interaction. We proceed by first finding the eigensystem of the operator  $(H_I(t) - i\frac{\partial}{\partial t})$  and treating  $H_t$  as a perturbation. The advantage of this approach is that the Floquet states are eigenstates of  $(H_I(t) - i\frac{\partial}{\partial t})$  and satisfy the following equation:  $[H_I(t) - i\frac{\partial}{\partial t}]\phi_\alpha(x, t) = \varepsilon_\alpha \phi_\alpha(x, t)$ . Hence, the corrections can be easily evaluated by working in an extended Hilbert space of  $T$ -periodic functions [24] by using

standard Rayleigh-Schrodinger perturbation theory. In our basis  $H_I(t)$  is diagonal and the resulting orthonormal set of eigenvectors for  $(H_I(t) - i\frac{\partial}{\partial t})$  are given by

$$|u_1(t)\rangle = \left( \exp\left[-i(U - \varepsilon - \varepsilon_1)t - i\frac{V}{\omega}\sin(\omega t)\right], 0, 0, 0 \right) \quad (21)$$

$$|u_2(t)\rangle = \left( 0, \exp\left[-i(U + \varepsilon - \varepsilon_2)t - i\frac{V}{\omega}\sin(\omega t)\right], 0, 0 \right) \quad (22)$$

$$|u_3(t)\rangle = \left( 0, 0, \exp(i\varepsilon_3 t), 0 \right) \quad (23)$$

$$|u_4(t)\rangle = \left( 0, 0, 0, \exp(i\varepsilon_4 t) \right) \quad (24)$$

Using the  $T$ -periodicity, the eigenvalues  $\varepsilon_{3,4}$  are zero and  $\varepsilon_1 = U + \varepsilon$  and  $\varepsilon_2 = U - \varepsilon$ . These eigenvalues represent zeroth order approximation to the quasi-energies. An interesting feature of these values is that for the choices  $U = \varepsilon$  and  $2U = n\omega$ , all four energies are degenerate. However, this degeneracy is lifted by  $H_t$  which is treated as a perturbation. The first order correction to the quasi-energies can be calculated in the extended Hilbert space of time periodic functions by defining an appropriate scalar product  $P_{ij} = \langle\langle u_i | H_t | u_j \rangle\rangle$ , where  $\langle\langle \dots \rangle\rangle$  denotes the inner product in the extended Hilbert space. It is straight forward to calculate matrix elements of  $P$  and subsequently its eigen values can be found. The matrix form of  $P$  is found to be:

$$\begin{pmatrix} 0 & 0 & \sqrt{2}W(-1)^n J_n(\beta) & -i\sqrt{2}\Omega(-1)^n J_n(\beta) \\ 0 & 0 & \sqrt{2}W J_n(\beta) & -i\sqrt{2}\Omega J_n(\beta) \\ \sqrt{2}W(-1)^n J_n(\beta) & \sqrt{2}W J_n(\beta) & V_+ & 0 \\ i\sqrt{2}\Omega(-1)^n J_n(\beta) & i\sqrt{2}\Omega J_n(\beta) & 0 & V_- \end{pmatrix} \quad (25)$$

where  $\beta = \frac{V}{\omega}$ . Thus we note that the electric field has redefined the tunneling and spin-orbit interaction in terms of the Bessel functions  $J_n(\beta)$ . The eigen values of the  $P$  give first order approximation to quasi-energies. Exact or near degeneracy of the quasi-energies results in the suppression of tunneling of the associated Floquet states, leading to localization of the initial state. This occurs whenever the ratio  $\frac{V}{\omega}$  is equal to the root of the Bessel function. Fig.3(b) shows the locations of the quasi-energies as a function of  $\frac{V}{\omega}$  from this perturbative treatment with  $n = 2$ .

Now we study the interplay between spin dynamics and the AC field we consider the time evolution of an initially localized state. We choose the system to be initially in the triplet state  $|T_0\rangle$ . The time evolution of an initial state is described by the Schrodinger equation. The wave function can be written as

$$|\Psi(t)\rangle = C_1(t)|(2, 0)S\rangle + C_2(t)|(0, 2)S\rangle + C_3(t)|(1, 1)S\rangle + C_4(t)|T_0\rangle \quad (26)$$

Focusing on  $H_1(t)$ , the evolution of the system is given by

$$i\dot{C}_1(t) = (U - \varepsilon - V\cos(\omega t))C_1(t) + \sqrt{2}WC_3(t) - i\sqrt{2}\Omega C_4(t) \quad (27)$$

$$i\dot{C}_2(t) = (U + \varepsilon + V\cos(\omega t))C_2(t) + \sqrt{2}WC_3(t) - i\sqrt{2}\Omega C_4(t) \quad (28)$$

$$i\dot{C}_3(t) = \sqrt{2}WC_1(t) + \sqrt{2}WC_2(t) + V_+ C_3(t) \quad (29)$$

$$i\dot{C}_4(t) = i\sqrt{2}\Omega C_1(t) + i\sqrt{2}\Omega C_2(t) + V_- C_4(t) \quad (30)$$

This system of equations can be solved numerically with the initial conditions  $(C_1(0), C_2(0), C_3(0), C_4(0) = 0, 0, 0, 1)$ . We will define the minimum value of  $|C_4(t)|^2$  in 20 driving periods as  $P_{(min)}^4$ . Hence the case  $P_{(min)}^4 = 1$  means that the system can maintain its initial state. In this case, tunneling between different states is completely suppressed. The case  $P_{(min)}^4 = 0$  means that the initial state cannot be maintained anymore. We also show time evolution of  $C_4(t)$  for different values of  $\frac{V}{\omega}$  and other parameters.

## VI. DISCUSSION

In this section we discuss our results. We first look at the behavior of  $P_{(min)}^4$  as a function of  $\frac{V}{\omega}$ , see Fig.4 We observe that in weak field regime  $P_{(min)}^4$  goes down as low as nearly zero. However as the electric field strength is

increased  $P_{(min)}^4$  forms peaks centered at specific value of  $\frac{V}{\omega}$  corresponding to close approach of two quasi-energies. This indicates the localization of initial state at these points. We can understand this behavior terms of behavior of quasi-energies. In Fig.3 we present our result of Floquet quasi-energy spectrum as a function of  $\frac{V}{\omega}$ . Fig.3(a) is obtained by numerical diagonalization of one period propagator  $U[T, 0]$  and in Fig.3(b) quasi-energies are obtained from Eqn.(25). The figure shows excellent agreement between perturbative result with ( $n = 2$ ) and the exact quasi-energies for strong and moderate electric fields. For weak fields, the terms in  $H_I(t)$  are not dominant over  $H_t$  terms and perturbation theory breaks down in weak field regime which explains the decay of  $P_{(min)}^4$  in this regime. The anti-crossings in quasi energies suppress quantum coherent tunneling and we observe spikes in  $P_{(min)}^4$  occur at these anti-crossings. Hence it can be concluded that the spikes in  $P_{(min)}^4$  given in Fig.3 correspond to localization of the initial state which occur at the anti-crossing points of quasi-energies.

Since our system is initially in the triplet state  $|T_0\rangle$ , we investigate the time evolution of this state through its amplitude  $C_4(t)$  for different values of the ration  $\frac{V}{\omega}$ , Fig.5. One can clearly see that  $C_4(t)$  remains unity for the entire time evolution for curve(a) which is plotted at the exact anti-crossing point. However, as we move away from the localization point,  $C_4(t)$  decays indicating that the initial state is no longer localized and tunneling is possible, see Fig.5 curves (b) and (c). Next we look at the effect of increasing spin-orbit field strength  $C_4(t)$ . Time evolution of the initial state is shown in Fig.6 at the anti-crossing point  $\frac{V}{\omega} = 14.79$  with spin-orbit strength set at 0.2. The curve is almost unity for entire evolution with a small oscillatory structure in  $C_4(t)$  between 1 and 0.8, indicating that this localization is robust against spin-orbit field strength. Thus the periodic field forms a well-established driven dynamics in the form of CDT even in the presence spin-flip mechanism. For comparison we consider time evolution of  $C_4(t)$  in the absence of AC field in Fig.7. We see that the state oscillates with spin-flip nature transitions as the triplet state can only couple to singlet states in our case. These oscillations are very rapid and do not have a strictly periodic structure. However the inclusion of the AC field prevents the spin from flipping and the state remains intact at as large values as  $\omega t = 400$ , manifesting CDT, see Fig.6.

For completeness we also discuss the behavior of the initial state for non-interacting electrons, given in Fig.8. In Fig. 8,  $\frac{V}{\omega}$  is set at different values but  $U = 0$ . The parameters are chosen as  $(\varepsilon, W, V_+, V_-, n) = (0, 0.1, 0, 0.05, 0)$  in units of  $\omega$ . Red curve is the 4th zero of Bessel function i.e 11.79. For black and orange this ratio is 11 and 12.5. We find that at exact anti-crossing,  $C_4(t)$  remains unity even for non interacting electrons. The state shows rapid oscillations at off resonance points as compared to interacting case shown in Fig.5.

A remark on the strength of the spin-orbit field  $\Omega$  is also due here. We refer to the values used in [9]. For GaAs dots the spin-orbit coupling  $\Omega$ , see Eqn.15, using  $W = 10\mu eV$ , an inter dot separation  $l = 50nm$ , and a spin-orbit length  $\lambda_{SO}$  in the range  $6 - 30\mu m$  we get  $|\Omega| = 20 - 110neV$ . However we have modeled spin-orbit field strength at higher value than this typical value. This can be achieved by adjusting inter dot separation  $l = 110nm$  [25] and tunneling strength  $W$ .

To summarize, we have shown that the DQD system with two interacting electrons under applied AC electric field can be used as a memory device. On the localization points, the initial triplet state maintains its memory and remains localized even though it is not the ground state of the system. Hence at these points spin is conserved even in the presence of spin-orbit interaction. Another interesting feature of this localization is that it is robust against changes in parameters like spin-orbit field strength and hopping amplitude in the system. It is expected that these results will be important in future spintronic device applications based on spins in DQD.

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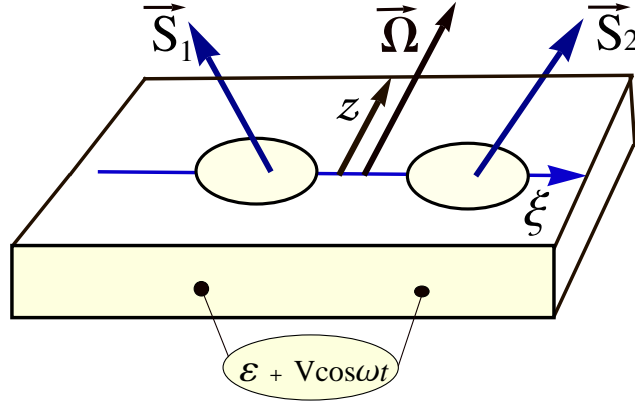


FIG. 1: The double quantum dot model and the coordinate system used:  $S_1$  and  $S_2$  denote the spin of the electron in the right and left quantum dots. The dots lie in  $(\xi, z)$  plane and tunneling is possible along  $\xi$  direction. The dots are detuned by externally applied voltage and an AC field is applied to the dots as shown.

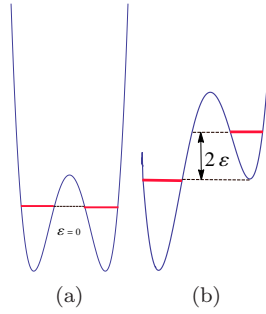


FIG. 2: (a) At zero detuning an electron will have same energy in both the dots. (b) For nonzero detuning an electron on the left dot is at  $2\epsilon$  lower energy than an electron on the right dot.

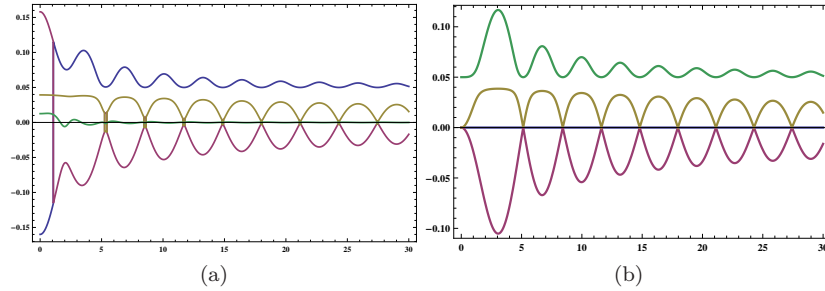


FIG. 3: Quasi-Energy spectrum as a function of  $\frac{V}{\omega}$ : Figure 3(a) is the perturbation theory result and 3(b) is the exact result. The parameters are chosen as  $(U, \epsilon, W, V_+, V_-, \Omega) = (1, 1, 0.1, 0, 0.05, 0.05)$  in units of  $\omega$ . In the figure quasi-energies and field strength  $V$  are measured in units of  $\omega$



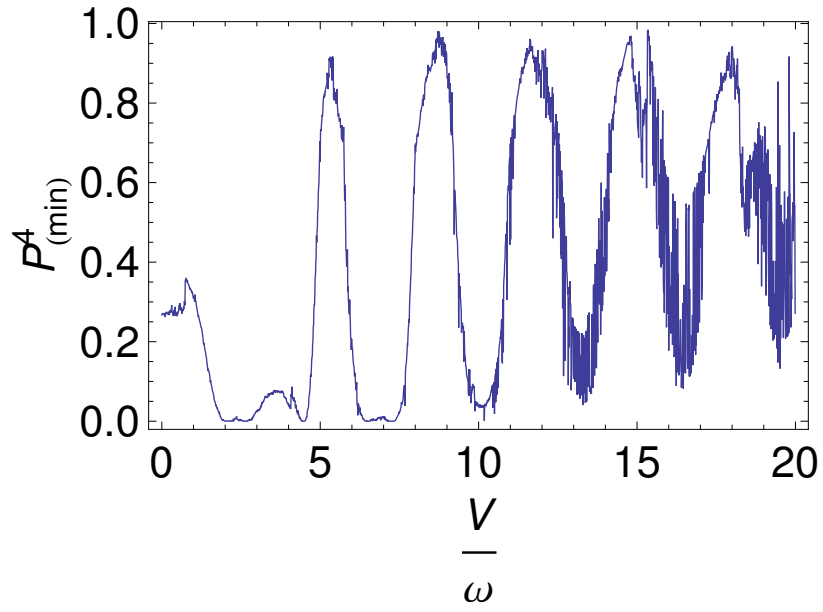


FIG. 4:  $P_{(min)}^4$  as a function of  $\frac{V}{\omega}$  indicating localization behavior of initial state

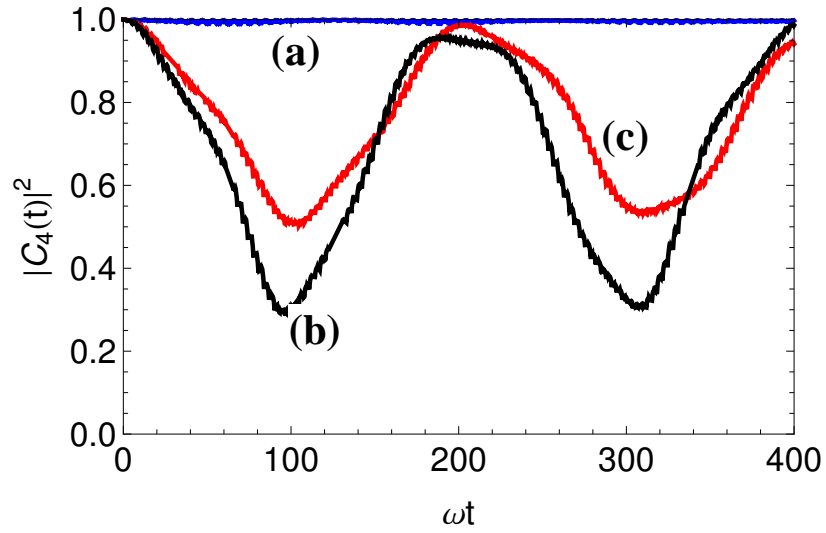


FIG. 5:  $C_4(t)$  as a function of  $\omega t$ : (a)  $\frac{V}{\omega} = 14.79$  the exact anti-crossing point. In this case  $C_4(t)$  remains unity for all times. Values of  $V$  are  $V = 13.5, 16$ , for (b),(c) respectively

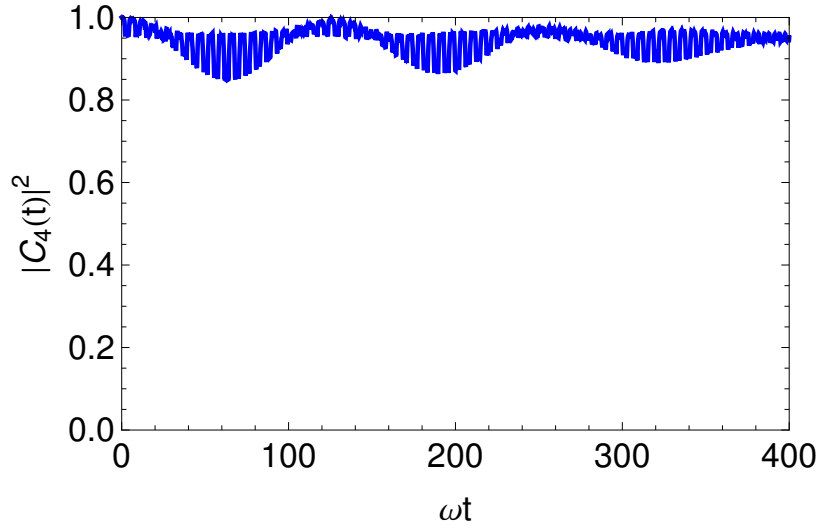


FIG. 6: The time evolution of the initial state with  $\Omega = 0.2$  at the anti-crossing point 14.79. Other parameters are  $(U, \varepsilon, W, V_+, V_-) = (1, 1, 0.1, 0, 0.05)$ .

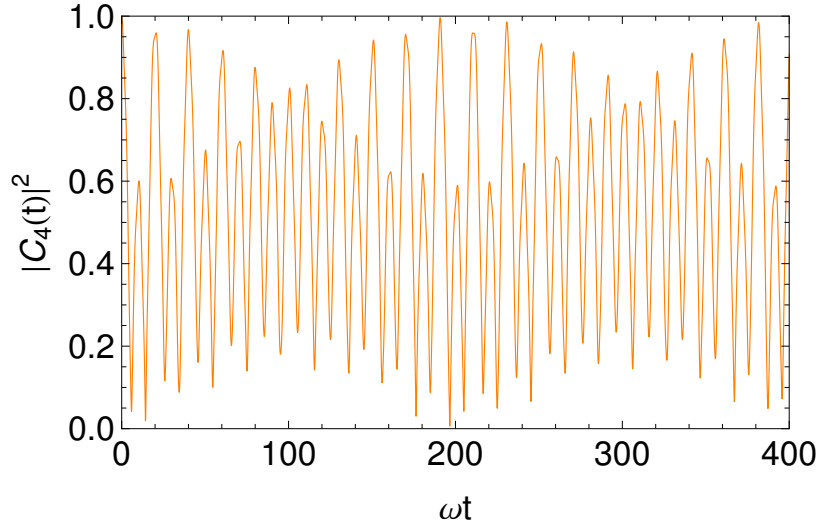


FIG. 7: The time evolution of  $C_4(t)$  in the absence of AC field. Parameters are the same as in Fig. 6

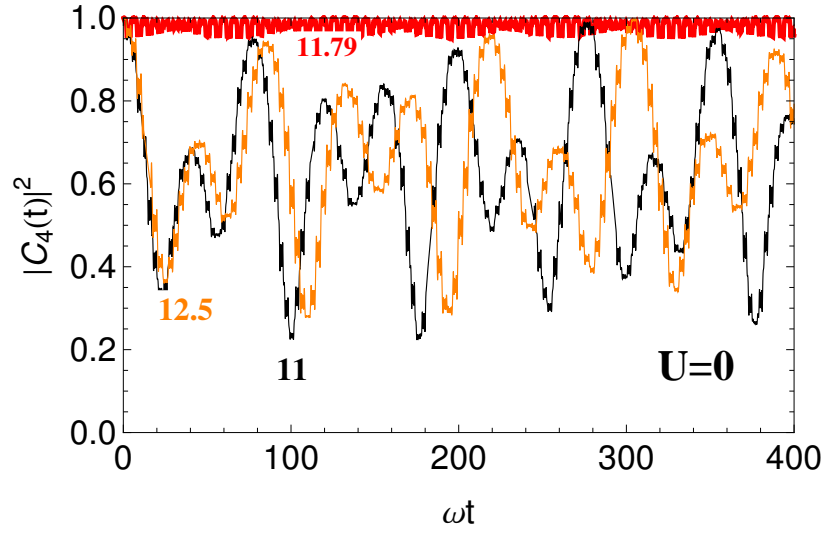


FIG. 8: Behavior of initial state for non interacting electrons for different values of  $\frac{V}{\omega}$  marked near the plots. Parameter chosen are  $(U, \varepsilon, W, V_+, V_-, \Omega) = (0, 0, 0.1, 0, 0.05, 0.2)$ . At exact anti-crossing  $C_4(t)$  remains unity even for non interacting case